SINP/TNP/01-10
hep-ph/0105057

Can R-parity violation lower $\sin 2\beta$?

Gautam Bhattacharyya [a,*], Amitava Datta [b,†] and Anirban Kundu [c,‡]

[a] Saha Institute of Nuclear Physics, 1/AF Bidhan Nagar, Kolkata 700064, India
[b] Department of Physics, Visva-Bharati, Santiniketan 731235, India
[c] Department of Physics, Jadavpur University, Kolkata 700032, India

Abstract

Recent time-dependent CP asymmetry measurements in the $B_d \to J/\psi K_S$ channel by the BaBar and Belle Collaborations yield somewhat lower values of $\sin 2\beta$ compared to the one obtained from the standard model fit. If the inconsistency between these numbers persists with more statistics, this will signal new physics contaminating the $B_d \to J/\psi K_S$ channel, thus disturbing the extraction of β. We show that the R-parity-violating interactions in supersymmetric theories can provide extra new phases which play a role in significantly reducing the above CP asymmetry, thus explaining why BaBar and Belle report lower values of $\sin 2\beta$. The same couplings also affect the $B_d \to \phi K_S$ decay rate and asymmetry, explain the $B \to \eta' K$ anomaly, and predict nonzero CP asymmetry in dominant B_s decays. The scenario will be tested in the ongoing and upcoming B factories.

PACS number(s): 11.30.Er, 13.25.Hw, 12.60.Jv, 11.30.Fs
Keywords: CP violation, B decays, supersymmetry, R-parity violation

Long after its discovery in the K system, evidence of CP violation is now being substantiated also in the B system, in particular, via the CP asymmetry measurement in the 'gold-plated' $B_d \to J/\psi K_S$ channel [1]. The CP asymmetry in the above channel is proportional to $\sin 2\beta$ in the standard model (SM), where $\beta = \text{Arg}\ (V_{td}^*)$ is an angle of the unitarity triangle of the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Any non-zero determination of β would be a signal of CP violation. The BaBar and Belle Collaborations, operating at the asymmetric B factories at SLAC and KEK respectively, have recently reported

$$\sin 2\beta = \begin{cases} 0.34 \pm 0.20 \pm 0.05 & \text{(BaBar [2])}, \\ 0.58^{+0.32+0.09}_{-0.34-0.10} & \text{(Belle [3])}. \end{cases} \tag{1}$$

When these are combined with the previous measurements, namely, $\sin 2\beta = 0.84 \pm 1.05$ (ALEPH [4]) and $\sin 2\beta = 0.79 \pm 0.44$ (CDF [5]), the global average reads

$$\sin 2\beta = 0.48 \pm 0.16. \tag{2}$$

On the other hand, using the experimental constraints from the measurement of $|\epsilon|$, $|V_{ub}/V_{cb}|$, Δm_d, and from the limit of Δm_s, the fitted value of $\sin 2\beta$, *strictly* within the framework of the SM, has been obtained as

$$\sin 2\beta = \begin{cases} 0.75 \pm 0.06\ [6], \\ 0.73 \pm 0.20\ [7]. \end{cases} \tag{3}$$

The two numbers in Eqs. (2) and (3) are at present consistent within errors, though their central values, as it should be noted, are fairly seperated. This separation may turn out to be an ideal new physics hunting ground. With an expected reduction of uncertainties of the parameters that go into the SM fit, and with an improved determination of the CP asymmetry in $B_d \to J/\psi K_S$ as the statistics accumulates[1], the

[*]Electronic address: gb@theory.saha.ernet.in
[†]Electronic address: adatta@juphys.ernet.in. On leave of absence from Jadavpur University, Kolkata 700032, India.
[‡]Electronic address: akundu@juphys.ernet.in
[1]The BaBar goal is to bring down the accuracy of $\sin 2\beta$ measurement to ± 0.06 with 30 fb^{-1} data [8].

inconsistency between Eqs. (2) and (3) may persist or may become even more prominant. In that case, a possible intervention of new physics with a new phase affecting the CP asymmetry in the $B_d \to J/\psi K_S$ channel cannot be ignored.

In the SM, the asymmetry in the $B_d \to J/\psi K_S$ channel is almost entirely mixing-induced. The time-dependent CP asymmetry is proportional to $\sin 2\beta$, where β appears in the phase of B_d-$\overline{B}_d$ mixing. The decay amplitude of $B_d \to J/\psi K_S$ does not carry any weak phase at leading order in the Wolfenstein parametrization. New physics might change the scenario in two ways. It can add a new weak phase in B_d-$\overline{B}_d$ mixing and/or it can generate a new diagram for $b \to c\bar{c}s$ decay amplitude that carries a new weak phase. Since the decay is Cabibbo-favoured, one usually tends to overlook the latter possibility. In this paper we examine the situation in which the decay is indeed affected by new physics, but the mixing amplitude is not. The point to note is that this scenario also induces a new physics amplitude in $B^+ \to J/\psi K^+$. Now two things can happen. First, direct CP asymmetry may be induced[2] in both $B_d \to J/\psi K_S$ and $B^+ \to J/\psi K^+$, which is non-existent in the SM. Indeed, for this to happen there must also exist a strong phase difference between the SM and new physics diagrams. Second, the mixing-induced CP asymmetry now depends not only on β, it involves a new weak phase as well. As a result, equating the CP asymmetry to $\sin 2\beta$ would be misleading. A combination of the 'true' β, the angle of the unitarity triangle, and other new parameters should now be related to the experimental CP asymmetry in the $B_d \to J/\psi K_S$ channel. This way it may be possible to explain why Eqs. (2) and (3) may disagree.

Since the values of $\sin 2\beta$ extracted from the CP asymmetry measurements in the asymmetric B factories have a tendency to be somewhat lower than the SM fit value, one is prompted to look for models where such tendency is favoured. In models of minimal flavour violation, where there are no new operators beyond those in the SM and no new weak phase beyond the one in the CKM matrix, a conservative scanning of all relevant input parameters in a standard CKM-like analysis yields an absolute lower bound on β. The SM, several versions of the minimal supersymmetric standard model (MSSM), and the two Higgs doublet models are examples of this class. The most conservative lower bound, as noted by the authors of Ref. [9, 10] is[3]

$$\sin(2\beta)_{\min} = \begin{cases} 0.42 \text{ (Present)}, \\ 0.48 \text{ (Future)}. \end{cases} \tag{4}$$

In particular, it has been shown in [9] that in the MSSM and in the minimal Supergravity (mSUGRA) models, the minimum $\sin 2\beta$ are

$$\sin(2\beta)_{\min} = \begin{cases} 0.40 \ (0.49) \text{ (MSSM)}, \\ 0.53 \ (0.62) \text{ (mSUGRA)}, \end{cases} \tag{5}$$

where the numbers within brackets correspond to future measurements of input parameters. In the context of supersymmetry, if CP violation in the K system is purely supersymmetric, i.e. ϵ and ϵ'/ϵ are completely explained by new phases of supersymmetric origin, then the CKM phase will be constrained from the charmless semileptonic B decays and B_d-$\overline{B}_d$ mixing. Only in that case, as the authors of Ref. [11] have argued, the CKM phase could be quite small leading to a very low $a_{\rm CP}$. The implications of a low $a_{\rm CP}$ in the context of a generic new physics scenario have been discussed in Refs. [12, 13].

The thrust of this paper is to examine the rôle of supersymmetry with broken R-parity [14] in the context outlined above. This brand of supersymmetry does not fall into the minimal flavour violating class, as it introduces new tree-level flavour changing operators. The essential points are outlined below. Recall that in the MSSM gauge invariance ensures neither the conservation of lepton number (L) nor that of baryon number (B). Using these quantum numbers R-parity is defined as $R = (-1)^{(3B+L+2S)}$, where S is the spin of the particle. R is $+1$ for all SM particles and -1 for their superpartners. In a general supersymmetric model one should in principle allow R-parity-violating ($\not{R}$) interactions. Tight constraints on the strength of these interactions exist in the literature [15]. Even though any concrete evidence for the existence of $\not{R}$ terms is still lacking, the observation of neutrino masses and mixings in solar and atmospheric neutrino data suggests that it would be premature to abandon the L-violating interactions

[2] An observation of direct CP violation even at a few percent level will constitute a definite signal of new physics.

[3] The bounds are conservative in the sense that they have been obtained by independently scanning all parameters under consideration.

[16]. Indeed, to avoid rapid proton decay one cannot simultaneously switch on both L- and B-violating interactions and for this reason we impose B conservation by hand. The L-violating superpotential that we consider here is $\lambda'_{ijk} L_i Q_j D^c_K$, where L_i and Q_i are lepton and quark doublet superfields, and D^c_i is down quark singlet superfields. The essential point is that the λ'-interactions can contribute to non-leptonic B decays at the tree level via slepton/sneutrino mediated graphs. In this paper we focus on $\not{R}$ effects arising at the $B_d \to J/\psi K_S$ decay amplitude level rather than through B_d-$\overline{B}_d$ mixing[4].

A detailed analysis of R-parity violation on $B_d \to J/\psi K_S$ and $B_d \to \phi K_S$ decay modes have earlier been carried out in Ref. [17]. We have extended the formalism of [17], which was based on mixing induced CP violation only, by incorporating the effects of the strong phase difference between the interfering amplitudes (an essential ingredient of direct CP violation). The latter constitutes the main source of CP violation in charged B decays and is present even in neutral B decays. The current availability of CLEO data on charged B decays together with the isospin symmetry enable us to extract quantitative results on the strong phase difference within this generalised framework. We also find that our essential conclusion regarding β extraction from $B_d \to J/\psi K_S$ channel is different from [17]; the reason is explained later.

In the SM, the matrix element of the effective Hamiltonian for $\overline{B}_d \to J/\psi K_S$ ($b \to c\bar{c}s$ at the quark level) is a combination of tree and penguin contributions, given by

$$\begin{aligned}
\langle J/\psi K_S | H_{\rm SM} | \overline{B_d} \rangle &= -\frac{G_F}{\sqrt{2}} \left(A^{\rm tree}_{\rm SM} + A^{\rm peng}_{\rm SM} \right), \quad \text{where,} \\
A^{\rm tree}_{\rm SM} &= V_{cb} V^*_{cs} (C_1 + \xi C_2) A_{J/\psi} \quad (\xi = 1/N_c), \\
A^{\rm peng}_{\rm SM} &= -V_{tb} V^*_{ts} (C_3 + \xi C_4 + C_5 + \xi C_6 + C_7 + \xi C_8 + C_9 + \xi C_{10}) A_{J/\psi}.
\end{aligned} \tag{6}$$

Here C_i's are the Wilson coefficients of the operators O_1-O_{10} defined as in [18] and evaluated at the factorization scale $\mu = m_b$, and

$$\begin{aligned}
A_{J/\psi} &= f_{J/\psi} F_0^{B\to K} f(m_B, m_K, m_{J/\psi}), \\
\text{with} \quad f(x,y,z) &= \sqrt{x^4 + y^4 + z^4 - 2x^2y^2 - 2z^2y^2 - 2x^2z^2}.
\end{aligned} \tag{7}$$

For numerical evaluation, we take the J/ψ decay constant $f_{J/\psi} = 0.38$ GeV [19] and the $B \to K$ decay form factor $F_0^{B\to K} = 0.42$ [20].

The λ'-induced interactions contribute to $b \to c\bar{c}s$ via slepton-mediated tree-level graphs. The matrix element of the effective $\not{R}$ Hamiltonian is given by

$$\langle J/\psi K_S | H_{\not{R}} | \overline{B_d} \rangle = -\frac{1}{4} u^R_{222} \, \xi \, A_{J/\psi}, \tag{8}$$

$$\text{where,} \quad u^R_{jnk} = \sum_{i=1}^{3} \frac{\lambda'^*_{in3} \lambda'_{ijk}}{2 m^2_{\tilde{e}_{iL}}}. \tag{9}$$

Due to the QCD dressing to the above operator, the expression in Eq. (8) should be multiplied by a factor ~ 2 at the scale m_b [21], which we have taken into account in all our numerical calculations.

The presence of $\not{R}$ terms modifies the expression for CP asymmetry in the $B_d \to J/\psi K_S$ channel in the following way. The key parameter of course is

$$\lambda_{J/\psi} = e^{-2i\beta} \frac{\langle J/\psi K_S | \overline{B}_d \rangle}{\langle J/\psi K_S | B_d \rangle} \equiv e^{-2i\beta} \frac{\overline{A}}{A}, \tag{10}$$

$$\text{where,} \quad A = A_{\rm SM}(1 + r e^{i\Delta\phi} e^{i\Delta\alpha}) \quad [\text{with } r = A_{\not{R}}/A_{\rm SM}], \tag{11}$$

$$\overline{A} = A_{\rm SM}(1 + r e^{-i\Delta\phi} e^{i\Delta\alpha}). \tag{12}$$

In the above expressions, $A_{\rm SM}$ and $A_{\not{R}}$ denote the signed magnitudes of the SM and $\not{R}$ amplitudes respectively[5], and we have ignored the overall phases in the expressions of A and $\overline{A}$. Here $\Delta\phi$ and $\Delta\alpha$ are

[4] Indeed, we respect the constraints on λ' couplings from Δm_d.

[5] $A_{\rm SM}$ contains both the tree and penguin amplitudes of the SM. Notice that the CKM factors in Eq. (6) are both real to a very good approximation. We also neglect the strong phase difference, expected to be small, that may appear between the tree and the $I = 1$ part of the electroweak penguin amplitudes. Among the different penguin contributions, only the dominant QCD part is included in our numerical estimates.

relative weak and strong phases between the SM and $\not{R}$ diagrams. The mixing induced CP asymmetry is given by

$$a_{\rm CP}^{\rm m} \equiv \frac{2{\rm Im}\lambda_{J/\psi}}{1+|\lambda_{J/\psi}|^2} = -\frac{2\rho}{1+\rho^2}\sin(2\beta+\zeta), \tag{13}$$

$$\text{where,}\ \ \rho = \frac{|\overline{A}|}{|A|} = \left(\frac{1+r^2+2r\cos(-\Delta\phi+\Delta\alpha)}{1+r^2+2r\cos(\Delta\phi+\Delta\alpha)}\right)^{1/2}, \tag{14}$$

$$\text{and,}\ \ \zeta = \tan^{-1}\left(\frac{2r\sin\Delta\phi(\cos\Delta\alpha + r\cos\Delta\phi)}{1+r^2\cos(2\Delta\phi)+2r\cos\Delta\phi\cos\Delta\alpha}\right). \tag{15}$$

The direct CP asymmetry is given by

$$a_{\rm CP}^{\rm d} \equiv \frac{1-|\lambda_{J/\psi}|^2}{1+|\lambda_{J/\psi}|^2} = -\frac{2r\sin\Delta\phi\sin\Delta\alpha}{1+r^2+2r\cos\Delta\phi\cos\Delta\alpha}. \tag{16}$$

The time-dependent CP asymmetry,

$$a_{\rm CP}(t) = \frac{B(B_d(t)\to J/\psi K_S) - B(\overline{B}_d(t)\to J/\psi K_S)}{B(B_d(t)\to J/\psi K_S) + B(\overline{B}_d(t)\to J/\psi K_S)}, \tag{17}$$

is given by

$$a_{\rm CP}(t) = a_{\rm CP}^{\rm d}\cos(\Delta m_d t) + a_{\rm CP}^{\rm m}\sin(\Delta m_d t). \tag{18}$$

After time integration, one obtains

$$a_{\rm CP} = \frac{1}{1+x^2}\left[a_{\rm CP}^{\rm d} + x a_{\rm CP}^{\rm m}\right];\ \ \text{where } x = (\Delta M/\Gamma)_{B_d,B_s}. \tag{19}$$

Side by side, a measurement of direct CP asymmetry in the $B^+ \to J/\psi K^+$ channel, which is the charged counterpart of $B_d \to J/\psi K_S$, yields important information about new physics. The asymmetry is defined by

$$a_{\rm CP}^{+} = \frac{B(B^+\to J/\psi K^+) - B(B^-\to J/\psi K^-)}{B(B^+\to J/\psi K^+) + B(B^-\to J/\psi K^-)}. \tag{20}$$

To a good approximation, $a_{\rm CP}^{+} = a_{\rm CP}^{\rm d}$ (see Ref. [22] for details). CLEO has measured [23]

$$a_{\rm CP}^{+} = (-1.8 \pm 4.3 \pm 0.4)\%. \tag{21}$$

The operators that mediate $B_d \to J/\psi K_S$ can have isospin I either 0 or 1. In fact, one can write the effective Hamiltonian in the $I=0$ and $I=1$ pieces in a model independent manner [22]. In the SM the $I=1$ contribution suffers a dynamical suppression. Recall that a sizable $a_{\rm CP}^{\rm d}$ necessarily requires a large strong phase difference (see Eq. (16)), which can result only from the interference between $I=0$ and $I=1$ amplitudes of comparable magnitude. In the SM, the former is far more dominant than the latter. As a result, $a_{\rm CP}^{\rm d}$ in the SM is vanishingly small. In some extensions beyond the SM, the $I=1$ piece may be slightly enhanced. A large enhancement however requires the presence of large rescattering effects, which is not a likely scenario [22]. In the present context, new physics contributes only to the $I=0$ sector by inducing a set of slepton mediated tree diagrams in the $b \to c\bar{c}s$ channel, for various combinations of λ' couplings. Thus in the absence of any possible enhancement of the $I=1$ part, the most likely scenario is that the strong phase difference ($\Delta\alpha$) is still vanishingly small, and so is the direct CP asymmetry. Yet, as we will see below, there is a significant impact of non-zero $\Delta\phi$ on the extraction and correct interpretation of β.

Following Eq. (14), ρ is identically equal to unity when $\Delta\alpha = 0$. Thus new physics can contaminate CP asymmetry only through

$$\zeta = \tan^{-1}\left(\frac{2r\sin\Delta\phi(1+r\cos\Delta\phi)}{1+r^2\cos(2\Delta\phi)+2r\cos\Delta\phi}\right). \tag{22}$$

To determine the maximum allowed size of r, we have to find out the experimental constraints on $\lambda'^{*}_{i23}\lambda'_{i22}$. The best constraints come from the measurements of $B^\pm \to \phi K^\pm$ branching ratio, which we will derive in this paper. The essential formalism [24] is described below. First note that due to the isospin structure of the interaction, the above product couplings contribute both to $b \to c\bar{c}s$ (i.e. $\overline{B}_d \to J/\psi K_S$ and $B^\pm \to J/\psi K^\pm$) and $b \to s\bar{s}s$ (i.e. $\overline{B}_d \to \phi K_S$ and $B^\pm \to \phi K^\pm$) at tree level. While for the former the diagram is slepton mediated, for the latter it is sneutrino mediated. For simplicity we assume that both the slepton and sneutrino are degenerate. Again note that for $B^\pm \to \phi K^\pm$ the leading SM diagram is a penguin, while the $\not{R}$ interaction, as mentioned before, proceeds at the tree level. The SM and $\not{R}$ effective Hamiltonians for $B^- \to \phi K^-$ lead to the following amplitudes:

$$\begin{aligned}\langle \phi K^- | H'_{\rm SM} | B^- \rangle &= \frac{G_F}{\sqrt{2}} V_{tb} V^*_{ts} \times \\ & \quad [(C_3 + C_4)(1+\xi) + C_5 + \xi C_6 - 0.5\{C_7 + \xi C_8 + (C_9 + C_{10})(1+\xi)\}]\, A_\phi, \quad (23)\\ \langle \phi K^- | H'_{\not{R}} | B^- \rangle &= -\frac{1}{4}\left(d^R_{222} + d^L_{222}\right)\xi\, A_\phi, \quad (24)\end{aligned}$$

where A_ϕ may be obtained analogously to the way $A_{J/\psi}$ was determined via Eq. (7), and

$$d^R_{jkn} = \sum_{i=1}^{3} \frac{\lambda'^{*}_{in3}\lambda'_{ijk}}{2m^2_{\tilde{\nu}_{iL}}}, \quad d^L_{jkn} = \sum_{i=1}^{3} \frac{\lambda'^{*}_{inj}\lambda'_{i3k}}{2m^2_{\tilde{\nu}_{iL}}}. \tag{25}$$

The most recent measurements of $B^\pm \to \phi K^\pm$ branching ratio are

$$B(B^\pm \to \phi K^\pm) = \begin{cases} \left(7.7^{+1.6}_{-1.4} \pm 0.8\right) \times 10^{-6} & \text{(BaBar [25])}, \\ \left(13.9^{+3.7+1.4}_{-3.3-2.4}\right) \times 10^{-6} & \text{(Belle [26])}, \\ \left(5.5^{+2.1}_{-1.8} \pm 0.6\right) \times 10^{-6} & \text{(CLEO [27])}. \end{cases} \tag{26}$$

Now for simplicity we assume that only one of d^R_{222} and d^L_{222} is non-zero, i.e. we do not admit unnatural cancellations between them. Since we are interested to put bounds on $\lambda'^{*}_{i23}\lambda'_{i22}$, we assume d^R_{222} to be non-zero. We further simplify the situation by assuming that only one combination, say the one corresponding to $i = 3$, is non-zero, i.e. the exchanged scalar is a tau-sneutrino. We also observe that the weak phase associated with this product coupling is totally arbitrary. In other words, $\lambda'^{*}_{323}\lambda'_{322}$ can take either sign, and $\Delta\phi$ is completely unconstrained. Note that the SM prediction for the $B^\pm \to \phi K^\pm$ varies in a wide range $(0.7 - 16) \times 10^{-6}$ (see Table I of Ref. [27]). To appreciate how much new physics effect we can accommodate, we assume that the SM contribution is close to the lower edge of the above range, and then saturate the 2σ CLEO upper limit in Eq. (26) *entirely* by $\not{R}$ interactions. This way we obtain[6]

$$|\lambda'^{*}_{323}\lambda'_{322}| \lesssim 1.5 \times 10^{-3}, \tag{27}$$

for an exchanged tau-sneutrino mass of 100 GeV. In fact, the above constraint is valid for any lepton family index i. This is the best constraint on the above combination[7], which we have derived for the first time in this paper.

To translate the limit in Eq. (27) into a limit on r, we need to decide in which regularization scheme $A_{\rm SM}$ will be computed. If the Wilson coefficients are computed in the 't Hooft-Veltman (HV) scheme[8], r lies in the range

$$-0.3 \lesssim r \lesssim 0.3. \tag{28}$$

[6]While deriving the bound in Eq. (27), we have multiplied the effective Hamiltonian in Eq. (24) by a QCD enhancement factor of 2, as was done with Eq. (8).

[7]Notice that the individual limits on λ'_{i23} and λ'_{i22} have been extracted from squark mediated processes assuming a mass of 100 GeV for whichever squark is involved [15]. While a 100 GeV slepton is very much consistent with all current data, the lower limit on a generic squark mass is presently pushed up to around 300 GeV from direct searches at Fermilab. Therefore, the λ' limits derived from squark mediated processes should be properly scaled while comparing them with those extracted from slepton exchanged diagrams.

[8]We have adapted the Wilson coefficients from Table 26 of Ref. [18]. The uncertainties in evaluating those coefficients arise from regularization scheme dependence, choice of $\Lambda^5_{\overline{\rm MS}}$, the factorization scale μ, and the long distance corrections. The Wilson coefficients we have used have been computed using $\Lambda^5_{\overline{\rm MS}} = 225$ MeV, $\mu = \overline{m}_b(m_b) = 4.4$ GeV and $m_t = 170$ GeV.

(There is no qualitative change in the result if we use other schemes like Naive Dimensional Regularization (NDR).) It is the negative value of r, which corresponds to a negative value of the $\not{R}$ product coupling, that has got an interesting implication in the extraction of β. As an illustration, putting $r = -0.3$ and $\Delta\phi = 90^\circ$ (say), it follows from Eq. (22) that $\zeta \sim -33^\circ$. Since our choices of λ' couplings do not contribute to Δm_d or any other observables that go into the SM fit, the latter still implies $\beta \sim 22^\circ$ (see Eq. (3)), which we accept as the 'true' value of β. As a result, the mixing induced CP asymmetry in the $B_d \to J/\psi K_S$ channel, as conceived via Eq. (13), now becomes ~ 0.2, as opposed to the SM expectation ~ 0.7. The crucial point is that a sizable negative ζ tends to cancel the 'true' 2β in the argument of the sine function in the expression of $a^{\rm m}_{\rm CP}$. This example demonstrates that $B_d \to J/\psi K_S$ need not be a 'gold-plated' channel for the determination of β. Rather, it could provide a window for new physics to manifest. Note that no drastic assumption, like large rescattering effects, etc, parametrized by a large $\sin\Delta\alpha$, was required to arrive at the above conclusion. In Figure 1, we demonstrate the variation of $\sin 2\beta$ with r for different values of $\Delta\phi$ and fixed $\Delta\alpha = 0$. The $\Delta\phi = 0$ curve corresponds to the SM reference value $\sin 2\beta = 0.7$. The minimum $\sin 2\beta$ we obtain in our scenario is

$$\sin(2\beta)_{\rm min} = 0.2, \tag{29}$$

which should be compared with Eqs. (4) and (5).

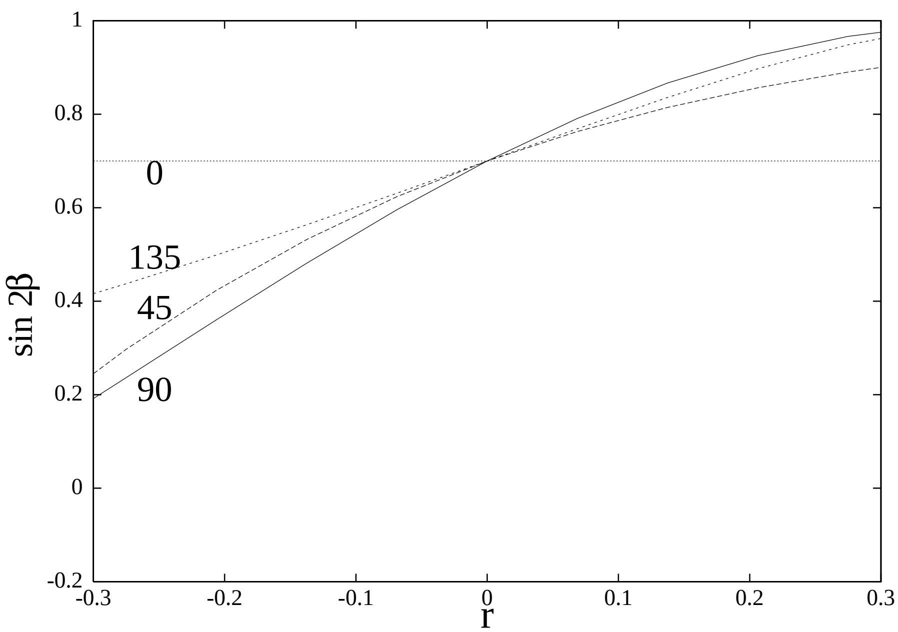

Figure 1: Variation of $(\sin 2\beta)_{J/\psi K_S}$ as a function of r and the weak phase difference $\Delta\phi$ (defined in the text). The values of $\Delta\phi$ (in degree) are indicated adjacent to the lines. The strong phase difference $\Delta\alpha$ is set to zero.

If we admit a fine-tuned situation $d^R_{222} = -d^L_{222}$ in Eq. (24), then the $B^\pm \to \phi K^\pm$ branching ratio constraints will not apply. But this product coupling will contribute to B_s-$\overline{B}_s$ mixing on which there is only an experimental lower limit. If we assume that in future the B_s-$\overline{B}_s$ mixing value settles, say, close to its present lower limit, then one requires $|\lambda'^*_{323}\lambda'_{322}| \sim 2.7 \times 10^{-3}$ (via sneutrino mediated box graphs [28]) to saturate the entire mixing by $\not{R}$ interaction. This is a conservative approach, as for a larger mixing one needs a larger value of the $\not{R}$ product coupling. With the above value of the product coupling, it is possible to arrange an even larger negative ζ, than obtained using Eq. (28), that can completely cancel the 'true' 2β inside the sine function of $a^{\rm m}_{\rm CP}$, which renders the latter almost zeroish. It must be admitted, though, that such a delicate fine-tuning is a very unlikely scenario.

So far we have worked putting $\Delta\alpha = 0$, i.e. by turning off any possible strong phase difference between the interfering amplitudes that might arise from long distance effects. Our approach has been a

conservative one. However, leaving aside any theoretical argument, one can put an experimental constraint on $\Delta\alpha$ for a given value of r and $\Delta\phi$. To do this we first take $r = -0.3$ and $\Delta\phi = 90^\circ$ as a reference point, which yielded the minimum $\sin 2\beta$ in our model. We then put these values in Eq. (16) and constrain $\Delta\alpha$ from the 2σ lower limit of the CLEO measurement of direct CP asymmetry in Eq. (21). We obtain $|\Delta\alpha| \lesssim 11^\circ$. This should be seen as an experimental constraint on long distance contributions for given values of the other parameters. Indeed, the BaBar and Belle collaborations can measure the sine and cosine time profiles of CP asymmetry in Eq. (18). The hadronic machines, on the other hand, are sensitive only to the time-integrated $a_{\rm CP}$ (see Eq. (19)). In the latter, the presence of a non-zero $a_{\rm CP}^{\rm d}$ can further disturb the extraction of β which resides inside $a_{\rm CP}^{\rm m}$.

The experimental branching ratio $B(B_d \to J/\psi K_S) = (9.5 \pm 1.0) \times 10^{-4}$ [29] is a factor of three to four higher than what one obtains from pure short distance effects involving naive factorization [30]. We do not attempt to solve this puzzle in this paper. We only make a remark that if the long distance contributions for short distance SM and $\not{R}$ operators behave similarly, then the long distance effects on $a_{\rm CP}$ will be minimal. Since the SM and $\not{R}$ operator structures are different, a separate study in this hitherto unexplored area is required. We do not get into those details in this paper.

A few comments on a previous analysis [17] of $\not{R}$ effects on $B_d \to J/\psi K_S$ are now in order. The bounds on $\not{R}$ couplings used in [17] are $|\lambda'^*_{i23}\lambda'_{i22}| \lesssim 1.4 \times 10^{-4}$ $(m_{\tilde{d}_R}/100 \text{ GeV})^2$ derived in [31]. As a result of such a strong constraint, the amplitude of $B_d \to J/\psi K_S$ changes only marginally. The above analysis led to a conclusion that the determination of β via $B_d \to J/\psi K_S$ was rather robust. In this paper, we counter this conclusion on the following ground. We notice that the bounds derived in [31] are 'basis-dependent', whose meaning is explained below. If only one λ' coupling in the weak basis is non-zero, more than one such coupling becomes non-zero in the mass basis, which are related to one another by the CKM elements. This way one can generate flavour-changing neutral currents (FCNCs) either in the down quark sector or in the up quark sector, depending on whether the CKM mixing is in the down sector or in the up sector respectively. Strong constraints on the λ' couplings emerge if one considers FCNCs completely in the down sector. If one chooses the other extreme, namely FCNCs entirely in the up sector, the constraints are not that tight. Since the choice of this basis is completely arbitrary, we prefer not to use such basis-dependent bounds in our analysis mainly because the conclusion depends heavily on this choice. Moreover, if one duly scales the exchanged down-squark mass to 300 GeV in view of its collider constraints, the limit used in [17] gets relaxed by an order of magnitude, and becomes closer to our limit in Eq. (27). A general study of CP violating B decays in nonleptonic modes in supersymmetric models with broken R-parity can also be found in Ref. [32].

The couplings $\lambda'^*_{i23}\lambda'_{i22}$ we have used in our analysis to reproduce a low $(\beta)_{J/\psi K_S}$ may have nontrivial impact on other processes as well. If a correlation is observed among the phenomena ocurring in these processes, it will certainly provide a strong motivation for the kind of new physics interactions we have advocated in our paper. These benchmark tests are listed below:

1. The SM amplitude of B_s-$\overline{B}_s$ mixing involves $V_{tb}V_{ts}^*$, and is real to a very good approximation. Hence we can expect the $b \to c\bar{c}s$ decays of the B_s meson (e.g. $B_s \to D_s^+ D_s^-$, $B_s \to J/\psi\phi$) to be CP conserving. But as mentioned before, the $\lambda'^*_{i23}\lambda'_{i22}$ couplings can contribute to B_s-$\overline{B}_s$ mixing through slepton mediated box graphs, which can interfere with the SM diagram. The same combinations, we have seen before, also affect the $b \to c\bar{c}s$ decay amplitude. The magnitude and the weak phase of the above product coupling needed to produce a low CP asymmetry in $B_d \to J/\psi K_S$ channel would necessarily give rise to a sizable CP asymmetry in the $B_s \to D_s^+ D_s^-$ and $B_s \to J/\psi\phi$ modes [33]. Note that neither R-parity conserving SUSY nor any other minimal flavour violating models can induce CP violation in the latter channels. In some sense, therefore, its observation will provide a necessary test for our scenario. The non-observation, on the other hand, will rule out our explanation of low $a_{\rm CP}$ in $B_d \to J/\psi K_S$. In any case, to observe CP violation in B_s decays, we have to wait till the second generation B factories, namely the LHC-b and BTeV, start taking data.

2. We have put constraints on the magnitude of the $\lambda'^*_{i23}\lambda'_{i22}$ couplings from the experimental $B^\pm \to \phi K^\pm$ branching ratio. In fact, the CP asymmetry in its neutral counterpart, namely the $B_d \to \phi K_S$ channel, is again expected to be proportional to $\sin 2\beta$ in the SM. Now since the SM amplitudes for $B_d \to J/\psi K_S$ and $B_d \to \phi K_S$ are quite different, it is quite apparent that in new physics inspired scenario, the CP

asymmetries measured in those two processes could yield values of β not only different from the 'true' β but also different from each other [17]. In other words, $\sin(2\beta)_{J/\psi K_S} \neq \sin(2\beta)_{\phi K_S} \neq \sin(2\beta)_{\rm SM\ fit}$. Again, to verify these non-equalities, we have to depend on the results from the second generation B factories.

3. It has been pointed out in Ref. [34] that the same $\lambda'^*_{i23}\lambda'_{i22}$ couplings can successfully explain the $B \to \eta' K$ anomaly [35]. It is noteworthy that the magnitude and phase of the above couplings required to produce a low $(\beta)_{J/\psi K_S}$ explains the anomaly by enhancing the SM branching ratio of $B \to \eta' K$ to its experimental value. Note again that none of the minimal flavour violating models can do this job.

4. It has been claimed that $\sin 2\beta$ can be determined very cleanly from the branching ratio measurements of the rare decays $K^+ \to \pi^+ \nu\overline{\nu}$ and $K^0_L \to \pi^0 \nu\overline{\nu}$ [36]. These processes, unlike those previously mentioned, will not be affected by our choice of $\not{R}$ couplings. Comparison of β extracted from these rare K decay processes with those obtained from several B decay channels may offer a powerful tool for probing physics beyond the SM.

To conclude, we have demonstrated that the distinctly lower central value of the BaBar measurement of CP asymmetry in the $B_d \to J/\psi K_S$ channel can be explained in models of supersymmetry with broken R-parity. In the process, we have derived new upper limits on the relevant $\not{R}$ couplings from the experimental $B^\pm \to \phi K^\pm$ branching ratio. It should be admitted though that the ability of the $\not{R}$ interactions to lower $(\beta)_{J/\psi K_S}$ is indeed shared by a few minimal flavour violating models, e.g. MSSM or mSUGRA. Thus if the disagreement between $(\beta)_{J/\psi K_S}$ and $(\beta)_{\rm SM\ fit}$ persists, this will signal new physics no doubt, but just with this single piece of information one cannot distinguish betwen the different models. What makes our scenario special is that it can do certain other jobs what the minimal flavour violating models cannot. We have outlined them above. Some of these tests can be carried out only in the second generation B factories. These tests can either boost our scenario or can rule it out.

Acknowledgments

The work of AD has been supported by the DST, India (Project No. SP/S2/k01/97) and the BRNS, India (Project No. 37/4/97 - R & D II/474). AK's work has been supported by the BRNS grant 2000/37/10/BRNS of DAE, Govt. of India, and by the grant F.10-14/2001 (SR-I) of UGC, India. We thank A. Raychaudhuri for reading the manuscript.

References

[1] A.B. Carter and A.I. Sanda, Phys. Rev. Lett. 45 (1980) 952; Phys. Rev. D 23 (1981) 1567; I.I. Bigi and A.I. Sanda, Nucl. Phys. B 193 (1981) 85.

[2] B. Aubert *et al.* (BaBar Collaboration), Phys. Rev. Lett. 86 (2001) 2515.

[3] A. Abashian *et al.* (Belle Collaboration), Phys. Rev. Lett. 86 (2001) 2509.

[4] R. Barate *et al.* (ALEPH Collaboration), Phys. Lett. B 492 (2000) 259; K. Ackerstaff *et al.* (OPAL Collaboration), Eur. Phys. J C5 (1998) 379.

[5] T. Affolder *et al.* (CDF Collaboration), Phys. Rev. D 61 (2000) 072005.

[6] F. Caravaglios, F. Parodi, P. Roudeau and A. Stocchi, hep-ph/0002171.

[7] A. Ali and D. London, hep-ph/0002167.

[8] The BaBar Physics Book: Physics at an asymmetric B factory; *eds:* P.F. Harrison and H.R. Quinn, SLAC Report SLAC-504, 1998.

[9] A. Buras and R. Buras, Phys. Lett. B 501 (2001) 223.

[10] A. Buras and R. Fleischer, hep-ph/0104238. See also, D. Atwood and A. Soni, hep-ph/0103197.

[11] A. Masiero and O. Vives, hep-ph/0104027; A. Masiero, M. Piai and O. Vives, hep-ph/0012096.

[12] G. Barenboim, G. Eyal and Y. Nir, Phys. Rev. Lett. 83 (1999) 4486.

[13] Y. Nir, hep-ph/0008226; G. Eyal, Y. Nir and G. Perez, JHEP 0008 (2000) 028; J.P. Silva and L. Wolfenstein, Phys. Rev. D 63 (2001) 056001; S. Bergmann and G. Perez, hep-ph/0103299.

[14] G. Farrar and P. Fayet, Phys. Lett. B 76 (1978) 575; S. Weinberg, Phys. Rev. D 26 (1982) 287; N. Sakai and T. Yanagida, Nucl. Phys. B 197 (1982) 533; C. Aulakh and R. Mohapatra, Phys. Lett. B 119 (1982) 136.

[15] For recent reviews, see G. Bhattacharyya, Nucl. Phys. B (Proc. Suppl.) 52A (1997) 83; hep-ph/9709395; H. Dreiner, hep-ph/9707435; R. Barbier *et al.*, hep-ph/9810232.

[16] R. Hempfling, Nucl. Phys. B 478 (1996) 3; H. Nilles and N. Polonsky, *ibid.* B 484 (1997) 33; E.J. Chun. S.K. Kang, C.W. Kim and U.W. Lee, Nucl. Phys. B 544 (1999) 89, E.J. Chun and J.S. Lee, Phys. Rev. D 60 (1999) 075006; S.Y. Choi, E.J. Chun, S.K. Kang and J.S. Lee, Phys. Rev. D 60 (1999) 075002; A. Joshipura and S. Vempati, Phys. Rev. D 60 (1999) 095009; S. Rakshit, G. Bhattacharyya and A. Raychaudhuri, Phys. Rev. D 59 (1999) 091701; G. Bhattacharyya, H.V. Klapdor-Kleingrothaus and H. Päs, Phys. Lett. B 463 (1999) 77.

[17] D. Guetta, Phys. Rev. D 58 (1998) 116008.

[18] A. Buras and R. Fleischer, in 'Heavy Flavours II', *eds*: A. Buras and M. Lindner (World Scientific, 1997).

[19] For a discussion, see M. Neubert and B. Stech, in 'Heavy Flavours II', *eds*: A. Buras and M. Lindner (World Scientific, 1997).

[20] M. Bauer, B. Stech and M. Wirbel, Z. Phys. C 34 (1987) 103.

[21] K. Huitu, C.-D. Lu, P. Singer and D.-X. Zhang, Phys. Rev. Lett. 81 (1998) 4313.

[22] R. Fleischer and T. Mannel, Phys. Lett. B 506 (2001) 311.

[23] G. Bonvicini *et al.* (CLEO Collaboration), Phys. Rev. Lett. 84 (2000) 5940.

[24] R. Fleischer and T. Mannel, hep-ph/0103121.

[25] B. Aubert *et al.* (BaBar Collaboration), hep-ex/0105001.

[26] A. Bozek (Belle Collaboration), talk at 4th International Conference on B Physics and CP violation, Ise, Japan, Feb 2001, hep-ex/0104041.

[27] R.A. Briere *et al.* (CLEO Collaboration), Phys. Rev. Lett. 86 (2001) 3718.

[28] G. Bhattacharyya and A. Raychaudhuri, Phys. Rev. D 57 (1998) 3837.

[29] Y. Gao, Talk given at DPF 2000 on behalf of the CLEO Collaboration, hep-ex/0011003.

[30] For a discussion, see A. Khodjamirian and R. Rückl, in 'Heavy Flavours II', *eds*: A. Buras and M. Lindner (World Scientific, 1997).

[31] K. Agashe and M. Graesser, Phys. Rev. D 54 (1996) 4445.

[32] J.-H. Jang and J.S. Lee, hep-ph/9808406.

[33] G. Bhattacharyya, A. Datta and A. Kundu, in preparation.

[34] D. Choudhury, B. Dutta and A. Kundu, Phys. Lett. B 456 (1999) 185.

[35] For the present status of the $B \to \eta' K$ anomaly, see S.J. Richichi *et al.* (CLEO Collaboration), Phys. Rev. Lett. 85 (2000) 520.

[36] G. Buchalla and A. Buras, Phys. Lett. B 333 (1994) 221; Phys. Rev. D 54 (1996) 6782.